

Heavy quarkonium production in double Pomeron exchange processes in perturbative QCD

Feng Yuan

Institut für Theoretische Physik, Universität Heidelberg, Philosophenweg 19, 69120 Heidelberg,

Germany

(March 2001)

Abstract

We calculate the heavy quarkonium production in double pomeron exchange processes in perturbative QCD by using two-gluon exchange model. For the P -wave χ_J productions, we find χ_1 and χ_2 production amplitudes which vanish in the forward scattering limit. We also calculate direct $J/\psi(\Upsilon) + \gamma$ production in the same approach, and these direct contributions are much smaller than the feeddown contributions from the P -wave states.

PACS number(s): 12.40.Nn, 13.85.Ni, 14.40.Gx

In recent years, there has been a renaissance of interest in diffractive scattering. These diffractive processes are described by the Regge theory in terms of the Pomeron (\mathcal{P}) exchange [1]. The Pomeron carries quantum numbers of the vacuum, so it is a colorless entity in QCD language, which may lead to the “rapidity gap” events in experiments. However, the nature of Pomeron and its reaction with hadrons remain a mystery. For a long time it had been understood that the dynamics of the “soft pomeron” is deeply tied to confinement [2]. However, it has been realized now that much can be learned about QCD from the wide variety of small- x and hard diffractive processes, which are now under study experimentally. In Refs. [3,4], the diffractive J/ψ and Υ production cross section have been formulated in photoproduction processes and in DIS processes in perturbative QCD. In the framework of perturbative QCD the Pomeron is represented by a pair of gluon in the color-singlet state. This two-gluon exchange model can successfully describe the experimental results from HERA [6].

This two-gluon exchange model has been extended to the hadron hadron collisions for the single diffractive processes [7,8], where the so-called coherent diffraction mechanism [9] plays an important role. In these processes, the Pomeron represented by a color-singlet two-gluon system is emitted from one hadron and interacts with another hadron in a hard process, in which the two gluons are both involved. These coherent diffraction processes are very interesting for the study of hard diffraction mechanism in hadron collisions, where there exist nonfactorization effects in the leading twist [9].

In this paper, we will further extend the above ideas to the double Pomeron exchange (DPE) processes. In particular, we will calculate heavy quarkonium productions in DPE processes at hadron colliders. Apart from the single diffractive processes, the DPE processes at $p\bar{p}$ or pp colliders are also useful to the studies of the diffraction mechanism and exploration of the nature of the Pomeron. In the DPE, both incoming hadrons are quasielastically scattered and therefore emerge at very large Feynman x_F , $|x_{F,1}|$, $|x_{F,2}| \geq 0.9$. And the final states system in the center region is produced from the Pomeron-Pomeron interaction. Due to the color-singlet (Pomeron) exchange, the final states system can be separated from the two beam particles (proton or antiproton) by large rapidity gaps. These DPE processes are quite attractive because it has been long time proposed to search Higgs meson in double rapidity gap events at proton colliders [10–16]. An obvious advantage of the rapidity gap approach is the spectacularly clean experimental signature [17].

In the two-gluon exchange model, χ_c and χ_b states productions in DPE processes have been estimated in Ref. [18], where they approximately take the gluons coupling to the χ states in $gg \rightarrow \chi_J$ from their decay widths in $\chi_J \rightarrow gg$ processes. In this paper, we will directly calculate these amplitudes in terms of the wave functions derivatives at the origin of the P -wave heavy quarkonium bound states in the framework of color-singlet model by using the nonrelativistic approximation, and we also take into account the strong correlations between the two incident gluons polarizations, which will enter into the production amplitudes explicitly. As a result, we find vanishing amplitudes for χ_1 and χ_2 states in the forward scattering limit. We will also calculate $J/\psi + \gamma$ production in DPE processes in the same approach.

The double pomeron exchange processes χ production in

$$p + p(\bar{p}) \rightarrow p + \chi_J + p(\bar{p}) \quad (1)$$

is plotted in Fig. 1 in the two gluon exchange model, where the hard subprocess $gg \rightarrow M$ is initiated by gluon-gluon fusion and the second t -channel gluon (with transverse momentum Q_T) is needed to screen the color flow across the rapidity gap intervals. In our calculations, we express the formulas in terms of the Sudakov variables. That is, every four-momenta k_i are decomposed as,

$$k_i = \alpha_i p_1 + \beta_i p_2 + \vec{k}_{iT}, \quad (2)$$

where p_1 and p_2 are the momenta of the incident hadrons (proton or antiproton). For high energy process, we will have $p_1^2 = 0$, $p_2^2 = 0$, and $2p_1 \cdot p_2 = s$. α_i and β_i are the momentum fractions of p_1 and p_2 respectively. \vec{k}_{iT} is the transverse momentum, which satisfies

$$\vec{k}_{iT} \cdot p_1 = 0, \quad \vec{k}_{iT} \cdot p_2 = 0. \quad (3)$$

For the momenta of the incident gluons q_1 and q_2 , we have the following decompositions,

$$q_1 = x_1 p_1 + q_{1T}, \quad q_2 = x_2 p_2 + q_{2T}. \quad (4)$$

And then, the momentum of the produced χ state will have the following form,

$$P(\chi) = q_1 + q_2 = x_1 p_1 + x_2 p_2 + P_T, \quad (5)$$

where the χ state transverse momentum $\vec{P}_T = \vec{q}_{1T} + \vec{q}_{2T}$. In the following, we will calculate the amplitudes in the forward scattering limit. That is to say, the momenta transfer squared between the final state hadrons and initiate hadrons are set to be zero, $t_1 = (p_1 - p'_1)^2 = 0$ and $t_2 = (p_2 - p'_2)^2 = 0$. Under this limit, we will have $\vec{q}_{1T} = -\vec{q}_{2T} = \vec{Q}_T$, and then $\vec{P}_T(\chi) = 0$. With the above kinematical considerations, we can formulate the forward scattering amplitudes for $pp \rightarrow p + \chi + p$ as follows,

$$\mathcal{M}(\chi_J) = \frac{9\pi^3}{4} \int \frac{d^2 Q_T}{\pi(Q_T^2)^2} f_g(x_1, x'_1, Q_T, M) f_g(x_2, x'_2, Q_T, M) \times A_J, \quad (6)$$

where f_g are the unintegrated off-forward (skewed) gluon distribution functions. They can be related to the conventional gluon density as [18]

$$f_g(x, x', Q_T, M) = R_g \frac{\partial}{\partial \ln Q_T^2} [\sqrt{T(Q_T, M)} x g(x, Q_T^2)]. \quad (7)$$

The factor R_g takes into account the skewed effects of the off-forward gluon density compared with the conventional gluon density in the region of $x' \ll x$. The bremsstrahlung survival probability T^2 can be found in [18], which will reduce to the conventional Sudakov form factors in the double logarithmic limit.

The expressions A_J in the amplitude (6) include the χ_J bound state effects (the wave functions), and will also depend on the polarizations of χ_J states. For $J = 0$, we have the following simple result,

$$A_0 = \frac{128\pi\alpha_s M}{9(M^2 + Q_T^2)^2} \sqrt{\frac{1}{4\pi M}} R'_P(0), \quad (8)$$

where $R'_P(0)$ is the radial wave function derivative at the origin for the P -wave χ state of charmonium or bottomonium. For $J = 1$,

$$A_1 = \frac{8 \times 128\pi\alpha_s}{9\sqrt{6}s(M^2 + Q_T^2)^2} \sqrt{\frac{1}{4\pi M}} R'_P(0) \epsilon_{\mu\nu\rho\sigma} e_{(J=1)}^\mu Q_T^\nu p_1^\rho p_2^\sigma, \quad (9)$$

where $e_{(J=1)}$ is the polarization vector of χ_1 state. From the above expression we can see that $A_1 \propto Q_T^\nu$, so the amplitude for $J = 1$ state will be zero after angular integration of (6) over \vec{Q}_T . That is to say, there is no contribution for χ_1 production in the DPE processes (1).

For $J = 2$,

$$A_2 = \frac{64\pi\alpha_s M}{9\sqrt{3}sQ_T^2(M^2 + Q_T^2)^2} \sqrt{\frac{1}{4\pi M}} R'_P(0) e_{(J=2)}^{\mu\nu} [4Q_T^2(p_{1\mu}p_{2\nu} + p_{1\nu}p_{2\mu}) + s(P_\mu P_\nu - 4Q_{T\mu}Q_{T\nu})]. \quad (10)$$

Here $e_{(J=2)}^{\mu\nu}$ is the polarization tensor for χ_2 state, which has the following properties,

$$e_{\mu\nu}P^\mu = 0, \quad e_{\mu\nu}g^{\mu\nu} = 0. \quad (11)$$

For the angular integration of (6) over \vec{Q}_T , we have the following identity,

$$\int d^2Q_T Q_T^\mu Q_T^\nu = \frac{\pi}{2} \int dQ_T^2 Q_T^2 g_{\mu\nu}^{(T)}, \quad (12)$$

where the transverse part of $g_{\mu\nu}$ is given by

$$g_{\mu\nu}^{(T)} = -g_{\mu\nu} + \frac{2}{s}(p_{1\mu}p_{2\nu} + p_{1\nu}p_{2\mu})$$

So, after integrating the azimuth angle of \vec{Q}_T , the expression in the brackets [] of Eq. (10) will become

$$[] = s(P_\mu P_\nu - 2g_{\mu\nu}Q_T^2). \quad (13)$$

With the above results, we find that the amplitude for χ_2 will be equal to zero because of the identities of Eq. (11).

The vanishing of χ_1 and χ_2 production amplitudes are direct consequences of the forward scattering limits in the DPE processes formulated in Eq. (6). Under this limit, the polarization vectors of the two incident gluons are strongly correlated, i.e., $\vec{e}_i \propto \vec{q}_{iT}$ with $\vec{q}_{1T} = -\vec{q}_{2T}$ [19–21,18]. Because of these strong correlations, the final state (χ_J state here) in the hard DPE processes must be in the $J_z = 0$ state [21]. On the other hand, we know that the 3P_2 state of $J_z = 0$ decouples with two real gluons in the nonrelativistic limit for the heavy quarkonium system [22]. By explicit calculations in the above, we show that the DPE amplitudes for χ_1 and χ_2 productions have the properties in (9) and (10), and finally give vanishing contribution to these states. We note that these properties are common results in the two-gluon exchange Pomeron model (either perturbative or nonperturbative two-gluon exchange models). However, in the other Pomeron model, such as D-L model [23], the production amplitudes for χ_1 and χ_2 do not vanish [24,25]. This is because in this model the

Pomeron couples approximately like a $C = +1$ photon, and the DPE χ_J amplitudes were adopted from the analogous process of $\chi_J \rightarrow \gamma^* \gamma^*$, which will give the same magnitude production rates for χ_1 and χ_2 as that of χ_0 [24,25].

Finally, the differential cross section of χ_0 production will be

$$d\sigma = \frac{|\mathcal{M}|^2}{16^2 \pi^3} e^{bt_1} e^{bt_2} dt_1 dt_2 dy, \quad (14)$$

where y is the rapidity of χ_0 state. t_i is the momentum transfer squared at the proton (antiproton) vertices, and b is slope for the proton form factor, which will be taken as $b = 4.0 \text{GeV}^{-2}$ in the following numerical calculations. After integrating t_1 and t_2 , we get

$$\frac{d\sigma}{dy} = \frac{|\mathcal{M}|^2}{16^2 \pi^3 b^2} = \frac{12\pi^4 \alpha_s^2 M}{b^2} |R'_P(0)|^2 \left[\int \frac{dQ_T^2}{(Q_T^2)^2} \frac{f_g(x_1, x'_1, Q_T, M) f_g(x_2, x'_2, Q_T, M)}{(M^2 + 2Q_T^2)^2} \right]^2. \quad (15)$$

In Fig. 2, we plot the differential cross sections (15) for χ_{c0} and $\chi_{b0}(1P)$ at the Fermilab Tevatron RUN II, where we adopt the rapidity gap survival probability to be $S^2 = 0.05$ as in Ref. [18]. For the masses and wave functions [26], we use

$$\begin{aligned} M(\chi_{c0}) &= 3.4 \text{GeV}, \quad |R'_P(0)|_c^2 = 0.075 \text{GeV}^5, \\ M(\chi_{b0}) &= 10.2 \text{GeV}, \quad |R'_P(0)|_b^2 = 1.42 \text{GeV}^5. \end{aligned}$$

For the gluon distribution function, we use the GRV LO parameterization [27]. After integrating over the rapidity, we get the following cross sections for double-diffractive productions of χ_{c0} and χ_{b0} as

$$\sigma(\chi_{c0}) = 735 \text{ nb}, \quad \sigma(\chi_{b0}) = 0.88 \text{ nb} \quad (16)$$

. Comparing our results with those of [18], we find that our result for χ_{c0} is roughly agreeing with [18]. However, our result for χ_{b0} is much larger than that in [18]. This is because we use the potential model calculations for the wave functions, while they use lattice calculations for the decay widths for χ_b states. So, for $b\bar{b}$ bound states, there exist large theoretical uncertainties associated with the bound states properties.

For the experimental observation, with the high resolution missing-mass measurements at the Tevatron ($\Delta M \approx 250 \text{MeV}$ [17]), we can identify all of the χ states in the DPE processes. Furthermore, we can also measure χ_0 productions in DPE by detecting their radiative decays to J/ψ (Υ) plus photon. The branching ratio of $\chi_0 \rightarrow J/\psi \gamma$ is about 6.6×10^{-3} , so the final cross section for J/ψ production from χ_{c0} feeddown decays is about $\sigma(\chi_{c0}) \times Br(\chi_{c0} \rightarrow J/\psi \gamma) \times Br(J/\psi \rightarrow \mu^+ \mu^-) \approx 300 \text{ pb}$. Having this large cross section, J/ψ production in the DPE processes can be well measured with large number of lepton pair ($\mu^+ \nu^-$ or $e^+ e^-$) events. For P -wave bottomonium states, we know that the branching ratio of $\chi_{b0}(1P) \rightarrow \Upsilon(1S)\gamma$ is very small (the upper bound is $BR < 6\%$), so it may be difficult to detect $\Upsilon(1S)$ from $\chi_{b0}(1P)$ decays. However, for $2P$ state $\chi_{b0}(2P)$, we have not so small branching ratios decaying into $\Upsilon(1S)$ and $\Upsilon(2S)$. So, from the combining branching ratios, we will have

$$\begin{aligned} \sigma(\chi_{b0}(2P)) \times & \left(Br(\chi_{b0}(2P) \rightarrow \Upsilon(2S)\gamma \rightarrow \mu^+ \mu^-) + Br(\chi_{b0}(2P) \rightarrow \Upsilon(1S)X \rightarrow \mu^+ \mu^-) \right) \\ & \approx 0.8 \text{ pb}. \end{aligned}$$

In the above estimate, we assumed that the production cross section of $\chi_{b0}(2P)$ is the same as that of $\chi_{b0}(1P)$ because their wave functions at the origin are roughly the same [26]. So, Υ states from the feeddown decays of $\chi_{b0}(2P)$ can also be observable in the DPE processes at the Tevatron II with integrated luminosity of 15fb^{-1} .

Apart from the feeddown contributions from χ_J states, J/ψ and Υ can also be produced directly from the DPE processes in

$$p + p(\bar{p}) \rightarrow p + J/\psi\gamma + p(\bar{p}), \quad (17)$$

for which we plot in Fig. 3 the typical diagram. For the final produced J/ψ , we will have the Sudakov variables as

$$P = \alpha_\psi p_1 + \beta_\psi p_2 + P_T,$$

and for the outgoing photon,

$$p_\gamma = \alpha_3 p_1 + \beta_3 p_2 + p_T^\gamma.$$

In the forward scattering limit, because $\vec{q}_{1T} + \vec{q}_{2T} = 0$, the transverse momentum of the photon will be equal to the transverse momentum of J/ψ in balance, i.e., $\vec{p}_T^\gamma + \vec{P}_T = 0$. The onshell conditions for the outgoing J/ψ and photon will give the following relations for the associated Sudakov variables,

$$\alpha_\psi \beta_\psi = \frac{P_T^2 + m^2}{s} \equiv \frac{m_T^2}{s}, \quad \alpha_3 \beta_3 = \frac{P_T^2}{s}. \quad (18)$$

From the momenta conservation, we also have

$$x_1 = \alpha_\psi + \alpha_3, \quad x_2 = \beta_\psi + \beta_3.$$

Using the above Sudakov variables, we derive the differential cross section for $J/\psi\gamma$ production in the DPE processes,

$$\frac{d\sigma}{dy_\gamma dy_\psi d^2 P_T} = \frac{|\mathcal{M}|^2}{16^3 \pi^6 b^2}, \quad (19)$$

where y_ψ and y_γ are the rapidities of final state J/ψ and photon, and α_ψ and α_3 can be related to rapidities,

$$\alpha_\psi = \frac{m_T}{\sqrt{s}} e^{y_\psi}, \quad \alpha_3 = \frac{P_T}{\sqrt{s}} e^{y_\gamma}.$$

The amplitude for $J/\psi + \gamma$ production is more complicated than those for χ_J productions. In the limit of $Q_T^2 \rightarrow 0$, we can simplify the amplitude and get the amplitude squared as

$$|\mathcal{M}|^2 = \frac{2 \times 16^3 \pi^8 \alpha_s^2 \alpha_e e_c^2 \alpha_3^2 \alpha_\psi^2 m |R_S(0)|^2}{3 x_1^2 (\alpha_3 m_T^2 + \alpha_\psi P_T^2)^2 (m_T^2)^2} \mathcal{I}_g^2, \quad (20)$$

where

$$\mathcal{I}_g = \int \frac{dQ_T^2}{(Q_T^2)^2} f_g(x_1, x'_1, Q_T, M) f_g(x_2, x'_2, Q_T, M). \quad (21)$$

So the cross section will be

$$\frac{d\sigma}{dy_\gamma dy_\psi d^2P_T} = \frac{2\pi^2 \alpha_s^2 \alpha_e e_c^2 \alpha_3^2 \alpha_\psi^2 m |R_S(0)|^2}{3b^2 x_1^2 (\alpha_3 m_T^2 + \alpha_\psi P_T^2)^2 (m_T^2)^2} \mathcal{I}_g^2. \quad (22)$$

In Fig. 4, we plot the differential cross sections $d\sigma/dy dP_T$ as functions of P_T at $y = 0$ for J/ψ and $\Upsilon(1S)$ at the Tevatron RUN II. For the numerical inputs for masses and wave functions, we use [26]

$$\begin{aligned} M(J/\psi) &= 3.1 \text{GeV}, \quad |R_S(0)|_c^2 = 0.81 \text{GeV}^5, \\ M(\Upsilon(1S)) &= 9.46 \text{GeV}, \quad |R_S(0)|_b^2 = 6.48 \text{GeV}^5. \end{aligned}$$

From Fig. 4, we see that the cross section for J/ψ drops very rapidly with P_T , and the cross section for Υ drops much more slowly than that of J/ψ . At low P_T , the cross section of J/ψ is much larger than that of Υ , but at large enough P_T it will be smaller than Υ . However, comparing the results in Fig. 4 to those in Fig. 2, we find that the cross sections for S -wave quarkonia are much smaller than those for P -wave states, even after the latter cross sections have been multiplied by the decay branching ratios of $\chi \rightarrow J/\psi(\Upsilon)X$. So, we conclude that the dominant contributions to S -wave quarkonium production in the DPE processes are the feeddown contributions from P -wave decays. This is quite different from the inclusive production processes or the DPE J/ψ production in the IS Pomeron model [28], where the direct production is the dominant source to the prompt J/ψ productions.

In conclusion, in this paper we have investigated the heavy quarkonium productions in the DPE processes at hadron colliders in perturbative QCD. We found that in the forward scattering limit, the amplitudes for χ_1 and χ_2 vanish due to the strong correlation between the polarization vectors of the two incident gluons. We also calculated the direct S -wave quarkonium production in the DPE processes associated with a photon, and found that the direct contributions are much smaller than the feeddown contributions from the P -wave states.

ACKNOWLEDGMENTS

We thank V.A. Khoze and M.G. Ryskin for correspondence and useful conversations and discussions. We are also grateful for the discussions with K.T. Chao, H.J. Pirner, and J.S. Xu.

REFERENCES

- [1] P.D.B. Collins, *An introduction to Regge theory and high energy physics*, Cambridge University Press, Cambridge (1977).
- [2] H. G. Dosch, T. Gousset, G. Kulzinger and H. J. Pirner, Phys. Rev. D **55**, 2602 (1997).
- [3] M.G. Ryskin, Z. Phys. **C37**, 89 (1993); M.G. Ryskin, et al., Z. Phys. **C76**, 231 (1996).
- [4] S.J. Brodsky et al., Phys. Rev. **D50**, 3134 (1994); L. Frankfurt et al., Phys. Rev. **D57**, 512 (1998).
- [5] T. Gehrmann and W.J. Stitling, Z. Phys. **C70**, 69 (1996); M. Genovese et al., Phys. Lett. **B378**, 347 (1996); E.M. Levin et al., hep-ph/9606443.
- [6] M. Derrick et al, ZEUS collaboration, Phys. Lett. **B350** 120 (1995); S. Aid et. al., H1 Collaboration, Nucl. Phys. B **472**, 3 (1996).
- [7] G. Alves, E. Levin, and A. Santoro, Phys. Rev. **D55**, 2683 (1997).
- [8] F. Yuan, J.S. Xu, H.A. Peng, and K.T. Chao, Phys. Rev. **D58**, 114016 (1998); F. Yuan and K.T. Chao, Phys. Rev. **D60**, 094009 (2000); **D60**, 094012 (2000); **D60**, 114018 (2000); **D60**, 117504 (2000).
- [9] J.C. Collins, L. Frankfurt and M. Strikman, Phys. Lett. **B307**, 161 (1993);
- [10] A. Schäfer, O. Nachtmann, and R. Schöpf, Phys. Lett. **B249**, 331 (1990).
- [11] B. Müller and A.J. Schramm, Nucl. Phys. **A523**, 677 (1991).
- [12] A. Bialas and P.V. Landshoff, Phys. Lett. **B256**, 540 (1991).
- [13] H.J. Lu and J. Milana, Phys. Rev. **D51**, 6107 (1995).
- [14] J.R. Cudell and O.F. Hernandez, Nucl. Phys. **B471**, 471 (1996).
- [15] V.A. Khoze, A.D. Martin, and M.G. Ryskin, Phys. Lett. **B401**, 330 (1997); Eur. Phys. J. **C14**, 525 (2000).
- [16] D. Kharzeev and E. Levin, hep-ph/0005311.
- [17] M.G. Albrow and A. Rostovtsev, hep-ph/0009336.
- [18] V.A. Khoze, A.D. Martin, and M.G. Ryskin, hep-ph/0011393.
- [19] J. Pumplin, Phys. Rev. **D52**, 1477 (1995).
- [20] A. Berera and J.C. Collins, Nucl. Phys. **B474**, 183 (1996).
- [21] V.A. Khoze, A.D. Martin, and M.G. Ryskin, Phys. Rev. **D56**, 5867 (1997); hep-ph/0006005; hep-ph/0010163; hep-ph/0103007.
- [22] A.I. Alekseev, Sov. Phys., JETP **34**, 826 (1958); R. Barbieri, R. Gatto, and R. Kögerler, Phys. Lett. **60B**, 183 (1976); M. Krammer and H. Krasemann, Phys. Lett. **73B**, 58 (1978).
- [23] A. Donnachie and P.V. Landshoff, Nucl. Phys. **B303**, 634 (1988); Phys. Lett. **B185**, 403 (1987); **B191**, 309 (1987).
- [24] E. Stein and A. Schäfer, Phys. Lett. **300**, 400 (1993).
- [25] H.A. Peng, et al., Phys. Lett. **B351**, 349 (1995); H.A. Peng and J.S. Xu, Commun. Theor. Phys. **29**, 287 (1998).
- [26] E.J. Eichten and C. Quigg, Phys. Rev. **D52**, 1726 (1995).
- [27] M. Glück et al. Z. Phys. **C67**, 433 (1995).
- [28] F. Yuan and K.T. Chao, J. Phys. **G25**, 1135 (1999).

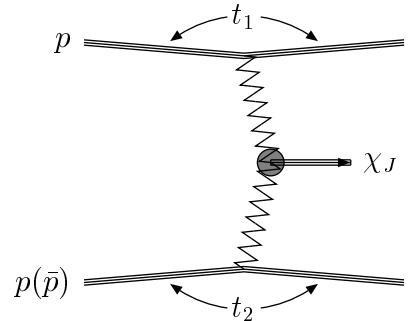
Figure Captions

Fig.1. Sketch diagram of the DPE processes χ_J production at hadron colliders: (a) Pomeron-Pomeron fusion into χ_J ; (b) Double Pomeron exchange in the perturbative two-gluon exchange picture.

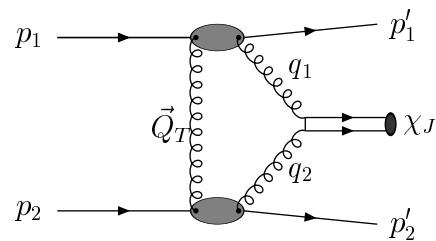
Fig.2. The differential cross sections $d\sigma/dy$ for χ_{c0} and χ_{b0} in DPE processes at the Fermilab Tevatron RUN II.

Fig.3. A typical diagram for the direct J/ψ (Υ) production in the DPE processes.

Fig.4. The differential cross sections $d\sigma/dydp_T$ for J/ψ and Υ production in the DPE at the Tevatron RUN II.



(a)



(b)

Fig.1 Sketch diagram of the DPE processes χ_J production at hadron colliders: (a) Pomeron-Pomeron fusion into χ_J ; (b) Double Pomeron exchange in the perturbative two-gluon exchange picture.

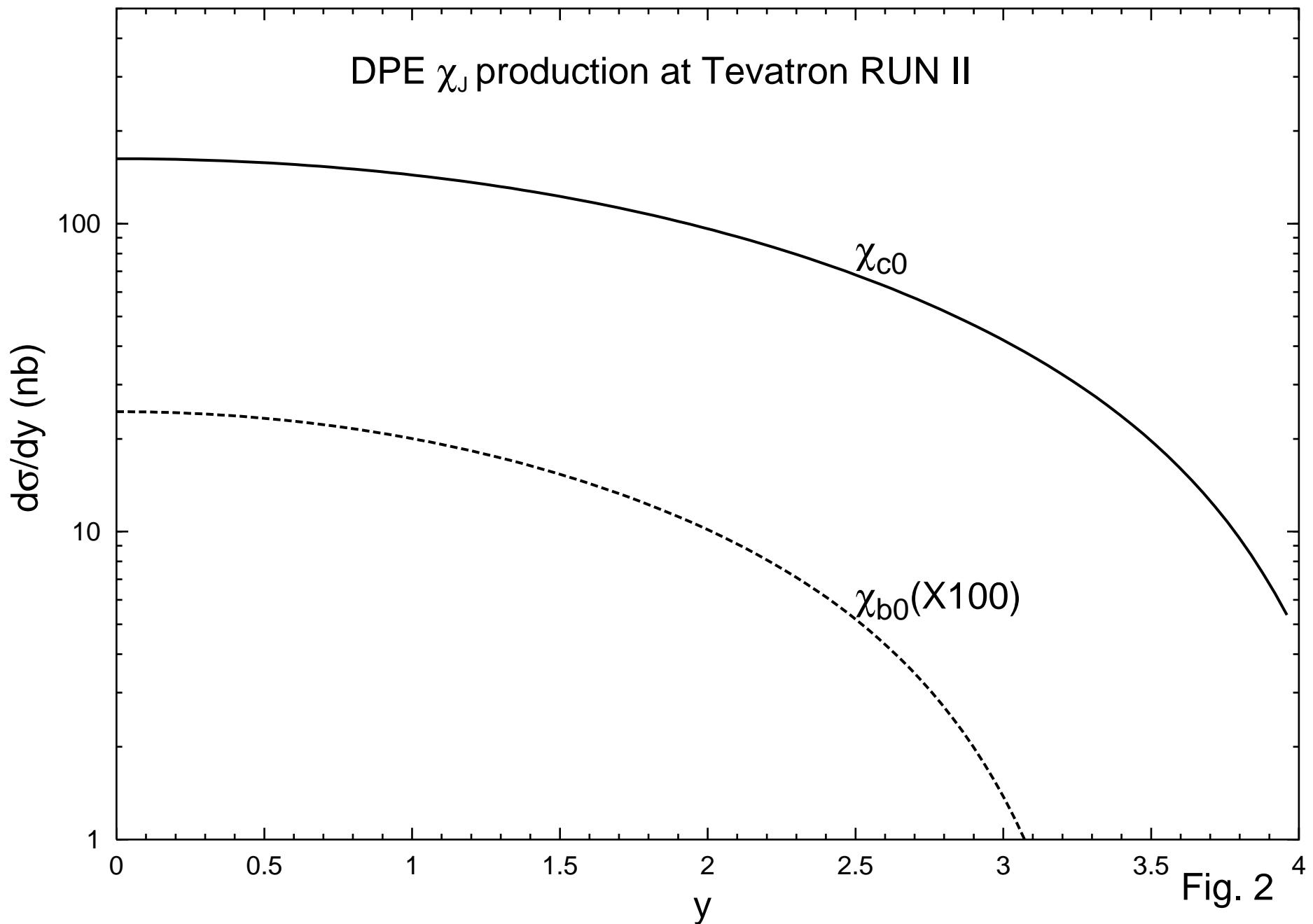


Fig. 2

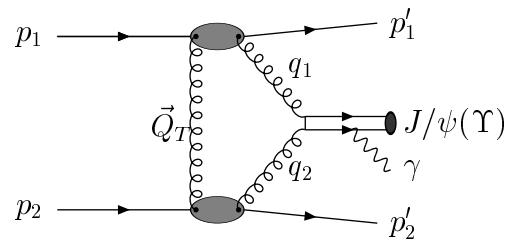


Fig.3

Fig.3 A typical diagram for the direct J/ψ (Υ) production in the DPE processes.

